

Coupling of $t\bar{t}$ and $\gamma\gamma$ with a strongly interacting Electroweak Symmetry Breaking Sector

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Abstract. We report the coupling of an external $\gamma\gamma$ or $t\bar{t}$ state to a strongly interacting EWSBS satisfying unitarity. We exploit perturbation theory for those coupling of the external state, whereas the EWSBS is taken as strongly interacting. We use a modified version of the IAM unitarization procedure to model such a strongly interacting regime. The matrix elements $V_L V_L \rightarrow V_L V_L$, $V_L V_L \leftrightarrow hh$, $hh \rightarrow hh$, $V_L V_L \leftrightarrow \{\gamma\gamma, t\bar{t}\}$, and $hh \leftrightarrow \{\gamma\gamma, t\bar{t}\}$ are all computed to NLO in perturbation theory with the Nonlinear Effective Field Theory of the EWSBS, within the Equivalence Theorem. This allows us to describe resonances of the electroweak sector that may be found at the LHC and their effect on other channels such as $\gamma\gamma$ or $t\bar{t}$ where they may be discovered.

1 Introduction

The new boson with a mass of 125 GeV found at the LHC [1, 2] would complete the Standard Model (SM) in its minimal version. The experimental collaborations at CERN are looking for deviations of its behaviour from that of the SM Higgs particle [3, 4]. The electroweak symmetry breaking happens at a scale of $v = 246$ GeV. New scalar resonances are constrained up to 600 – 700 GeV [5]. The constraint over new vector bosons reaches even higher in energy [6]. Thus, the study of the TeV scale is the goal of LHC Run-II. And it is expectable that new physics related with the Electroweak Symmetry Breaking Sector (EWSBS), if it exists, will be found at this scale.

To model the behaviour of a generic EWSBS with a theory which only includes the particles already known to exist there (the new Higgslike boson h and the longitudinal components of gauge bosons W_L^\pm, Z_L), the so-called Higgs Effective Field Theory (HEFT) is used [7–22]. A summary of these efforts can be found on the 4th CERN Yellow Report [23]. This model is based on the old (Higgsless) Electroweak Chiral Lagrangian (ECL) [24], which is also inspired by Chiral Perturbation Theory (ChPt) for hadron physics [25].

In our Refs. [26–30], we simplify the study of the HEFT by means of the Equivalence Theorem [31], which is valid for

$$s \gg M_h^2, M_W^2, M_Z^2 \sim (100 \text{ GeV})^2. \quad (1)$$

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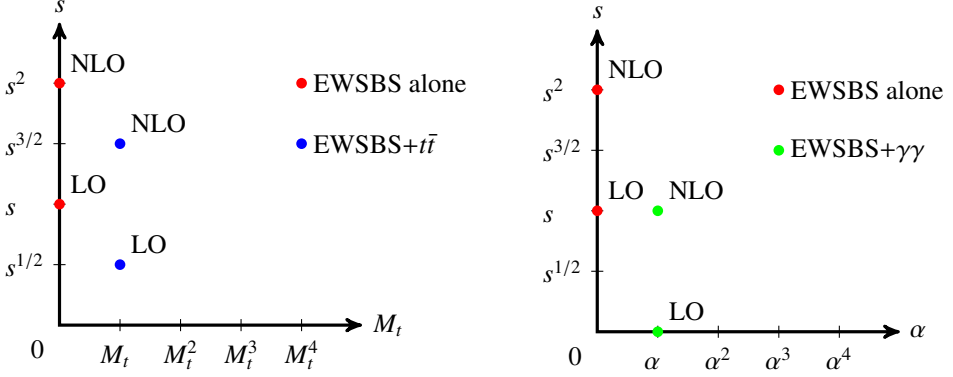


Figure 1. Left: chiral (M_t, \sqrt{s}) counting for coupling with $t\bar{t}$. Right: chiral (α, s) counting for coupling with $\gamma\gamma$. Note that the HEFT is perturbative in M_t/v and α_{EM} , respectively, but requires unitarization in $s/(4\pi v)^2$ (all orders along the OY axis included) to reach the resonance region.

Since the possible resonance region is above 500 GeV, this approximation is safe, and allows us to identify the longitudinal W_L^\pm, Z_L with the pseudo-Goldstone bosons of symmetry breaking ω^a ($a = 1, 2, 3$) in their scattering amplitudes. For instance,

$$T(W_L^i W_L^j \rightarrow W_L^k W_L^l) = T(\omega^i \omega^j \rightarrow \omega^k \omega^l) + \mathcal{O}\left(\frac{M_W}{\sqrt{s}}\right). \quad (2)$$

We have found [26–30] that, for any parameter choice separating from the SM, the theory becomes strongly interacting at sufficiently high energy, and resonances may appear. Dispersion relations (as unitarization procedures) are used to deal with this non-perturbative regime. This also happens for the ChPT theory for hadron physics [32, 33]. In Ref. [34, 35], unitarization procedures are tested for the hadron case of ChPT.

In this work, based on our recent works [36, 37], we couple $\gamma\gamma$ and $t\bar{t}$ states to the unitarized scattering amplitudes of our refs. [28, 30]. The couplings between the EWSBS and $\{t\bar{t}, \gamma\gamma\}$ states are perturbative in M_t/v (negligible if compared with \sqrt{s}/v) and α_{EM} , respectively. However, the EWSBS itself is strongly interacting (see Fig. 1). For the $t\bar{t}$ case, account must be taken that we are in the regime $M_t^2/v^2 \ll \sqrt{s}M_t/v^2 \ll s/v^2$. This allows us to neglect diagrams like the 1-loop top mass renormalization, since they contribute to the scattering amplitudes with higher orders of \sqrt{s}/v . Note also that, in the high energy limit $s \ll M_{Z,W}^2 \sim M_h^2$ where the Equivalence Theorem can be applied, we can take $M_W^2 = M_Z^2 = M_h^2 = 0$ consistently.

Note that we compute $\{\omega, \omega\} \leftrightarrow \{\omega\omega, hh, \gamma\gamma, t\bar{t}\}$. $\gamma\gamma$ and $t\bar{t}$ can both appear as initial and final states. Of course, $\gamma\gamma$ and $t\bar{t}$ as final states are being tested at the LHC experiments, and physics of two photon decays has been pursued since the dawn of particle physics [38]. But we should not forget the physics of $\gamma\gamma \rightarrow \omega\omega$, since the photon can be a parton of the proton [39] or the electron in pp and e^-e^+ colliders, respectively. The CMS collaboration [40] is currently setting bounds to anomalous quartic gauge couplings by analysing $\gamma\gamma \rightarrow W^+W^-$. Photon colliders driven by lepton beams due to Compton backscattering could also become a future application of this work [41, 42].

2 The Electroweak Chiral Lagrangian

First we quote the HEFT Lagrangian [29, 30, 37] which includes only the low-energy dynamics of the quark sector and the four light modes: three would-be Goldstone Bosons ω^a (WBGBs) and the Higgs-like particle h . A global symmetry breaking pattern $SU(2)_L \times SU(2)_R \rightarrow SU(2)_C$ is considered.

$$\mathcal{L} = \frac{v^2}{4} \mathcal{F}(h/v) \text{Tr}[(D_\mu U)^\dagger D^\mu U] + \frac{1}{2} \partial_\mu h \partial^\mu h - V(h) + i \bar{Q} \partial Q - v \mathcal{G}(h/v) [\bar{Q}'_L U H_Q Q'_R + \text{h.c.}]. \quad (3)$$

The $U(x) \in SU(2)_C$ isospin can be parametrized by using the so-called spherical parametrization [30] by means of

$$U = \sqrt{1 - \frac{\omega^2}{v^2}} + i \frac{\omega^i}{v} \tau_i. \quad (4)$$

Note that $i = 1, 2, 3$ in the isospin basis, which is related to the charge one by $\omega^\pm = (\omega^1 \mp i\omega^2)/\sqrt{2}$ and $\omega^0 = \omega^3$. The $SU(2)_L \times U(1)_Y$ covariant derivative is given by

$$D_\mu U = \partial_\mu U + i \frac{g}{2} \tau_i U W_\mu^i - i \frac{g'}{2} U \tau_3 B_\mu = i \frac{\partial_\mu \omega^i}{v} \tau_i + i \frac{g}{2} W_\mu^i \tau_i - i \frac{g'}{2} B_\mu \tau_3 + \dots, \quad (5)$$

where the dots represent terms of higher order in (ω^a/v) .

The Higgs potential is expanded as

$$V(h/v) = v^4 \sum_{n=3}^{\infty} V_n \left(\frac{h}{v} \right)^n. \quad (6)$$

The SM is recovered for $V_3 = M_h^2/2v^2$, $V_4 = M_h^2/8v^2$, $V_n = 0 \ \forall n > 4$. Note that these terms are subleading in our approximation (Eq. 1), so that we could neglect the whole potential V provided that this behaviour holds by whatever beyond SM theory happens to succeed. This is the case in most models of interest, and is a reasonable hypothesis since the constraints of these couplings have so far been found to be close to the SM values.

By expanding Eq. (3) without considering (yet) the Yukawa part, we obtain

$$\begin{aligned} \mathcal{L} = & \frac{1}{2} \partial_\mu h \partial^\mu h + \frac{1}{2} \mathcal{F}(h/v) (2 \partial_\mu \omega^+ \partial^\mu \omega^- + \partial_\mu \omega^0 \partial^\mu \omega^0) \\ & + \frac{1}{2v^2} \mathcal{F}(h/v) (\partial_\mu \omega^+ \omega^- + \omega^+ \partial_\mu \omega^- + \omega^0 \partial_\mu \omega^0)^2 \\ & + ie \mathcal{F}(h/v) A^\mu (\partial_\mu \omega^+ \omega^- - \omega^+ \partial_\mu \omega^-) + e^2 \mathcal{F}(h/v) A_\mu A^\mu \omega^+ \omega^-, \end{aligned} \quad (7)$$

where $A_\mu = \sin \theta_W W_{\mu,3} + \cos \theta_W B_\mu$. We are not considering couplings with external transverse gauge bosons. Note that these neglected states do not appear in inner loops since this would be a higher order correction in α_{EM} .

The Yukawa part of the Lagrangian of Eq. (3), once the Yukawa-coupling matrix is diagonalized [37], can be written as

$$\begin{aligned} \mathcal{L}_Y = & -\mathcal{G}(h) \left[\sqrt{1 - \frac{\omega^2}{v^2}} (M_t \bar{t} t + M_b \bar{b} b) + \frac{i\omega^0}{v} (M_t \bar{t} \gamma^5 t + M_b \bar{b} \gamma^5 b) \right. \\ & \left. + i\sqrt{2} \frac{\omega^+}{v} (M_b \bar{t}_L b_R - M_t \bar{t}_R b_L) + i\sqrt{2} \frac{\omega^-}{v} (M_t \bar{b}_L t_R - M_b \bar{b}_R t_L) \right]. \end{aligned} \quad (8)$$

Note that we have considered only couplings with the third quark generation ($M_{t,b} \gg M_{c,s,u,d} \rightarrow 0$). This Lagrangian breaks custodial symmetry because of $M_t \gg M_b$.

The \mathcal{F} and \mathcal{G} functions of Eq. (3) are parametrized as

$$\mathcal{F}(h/v) = 1 + 2a \frac{h}{v} + b \frac{h^2}{v^2} + \dots \quad \mathcal{G}(h/v) = 1 + c_1 \frac{h}{v} + c_2 \frac{h^2}{v^2} + \dots, \quad (9)$$

so that Eq. (8) can be written as

$$\mathcal{L}_Y = - \left(1 + c_1 \frac{h}{v} + c_2 \frac{h^2}{v^2} \right) \left[\left(1 - \frac{\omega^2}{2v^2} \right) M_t \bar{t} t + \frac{i\omega^0}{v} M_t \bar{t} \gamma^5 t - i\sqrt{2} \frac{\omega^+}{v} M_t \bar{t}_R b_L + i\sqrt{2} \frac{\omega^-}{v} M_t \bar{b}_L t_R \right]. \quad (10)$$

Since this is an EFT, in order to renormalize the tree level Lagrangian in Eq. (7) at the one-loop level, counterterms of dim. 8 are needed. For the EWSBS (strong) interactions, the minimal set of counterterms is $\{a_4, a_5, d, e, g\}$ (see Refs. [27, 30]),

$$\begin{aligned} \mathcal{L}_{4,\text{EWSBS}} = & \frac{4a_4}{v^4} \partial_\mu \omega^i \partial_\nu \omega^j \partial^\mu \omega^j \partial^\nu \omega^j + \frac{4a_5}{v^4} \partial_\mu \omega^i \partial^\mu \omega^j \partial_\nu \omega^j \partial^\nu \omega^j \\ & + \frac{2d}{v^4} \partial_\mu h \partial^\mu h \partial_\nu \omega^i \partial^\nu \omega^i + \frac{2e}{v^4} \partial_\mu h \partial^\nu h \partial^\mu \omega^i \partial_\nu \omega^i + \frac{g}{v^4} (\partial_\mu h \partial^\mu h)^2 \end{aligned} \quad (11)$$

For the $\gamma\gamma$ coupling [29, 36] we consider $\{a_1, a_2, a_3, c_\gamma\}$, though only $a_1 + a_2 - a_3$ appears and none is strictly needed,

$$\begin{aligned} \mathcal{L}_{4,\gamma\gamma} = & \frac{e^2 a_1}{2v^2} A_{\mu\nu} A^{\mu\nu} (v^2 - 4\omega^+ \omega^-) + \frac{2e(a_2 - a_3)}{v^2} A_{\mu\nu} [i(\partial^\nu \omega^+ \partial^\mu \omega^- - \partial^\mu \omega^+ \partial^\nu \omega^-) \\ & + eA^\mu (\omega^+ \partial^\nu \omega^- + \omega^- \partial^\nu \omega^+) - eA^\nu (\omega^+ \partial^\mu \omega^- + \omega^- \partial^\mu \omega^+)] - \frac{c_\gamma}{2} \frac{h}{v} e^2 A_{\mu\nu} A^{\mu\nu}. \end{aligned} \quad (12)$$

And for the Yukawa coupling with top quarks [37],

$$\mathcal{L}_{4,\bar{t}t} = g_t \frac{M_t}{v^4} (\partial_\mu \omega^i \partial^\mu \omega^j) \bar{t} t + g'_t \frac{M_t}{v^4} (\partial_\mu h \partial^\mu h) \bar{t} t \quad (13)$$

3 Partial waves

The perturbative $\{\omega^j \omega^j, hh\} \rightarrow \{\omega^k \omega^l, hh\}$ scattering amplitudes can be found in our Ref. [30]. For the representation of the $\omega\omega$ states, we will use the isospin basis $|I, M_I\rangle$ (I is the isospin and M_I , its projection). Because of isospin symmetry, scattering amplitudes factorizes in this basis, and their values do not depend on M_I . Unitarization procedures (based on dispersion relations) are most easily applicable over a partial wave decomposition [30] which, for $\{\omega^j \omega^j, hh\}$ states, are computed by means of

$$A_{IJ}(s) = \frac{1}{64\pi} \int_{-1}^1 d(\cos \theta) P_J(\cos \theta) A_I(s, t, u). \quad (14)$$

Since the hh state is an isospin singlet $I = 0$ (which couples with $J = 0, 2$), if $I \neq 0$ there is no mixing with hh channel, and the non-vanishing matrix elements are [30] $IJ = 11, 20, 22$. A similar expression to Eq. (14) is used for $M_I(\omega^i \omega^j \rightarrow hh)$ and $T_I(hh \rightarrow hh)$ ($J = 0$ in these case).

Note that partial waves of Eq. (14) have a chiral expansion

$$A_{IJ}(s) = A_{IJ}^{(0)}(s) + A_{IJ}^{(1)}(s) + \dots, \quad (15)$$

where $A_{IJ}^{(0)}(s) \sim s$ corresponds to the LO term and $A_{IJ}^{(1)}(s) \sim s^2$, to the NLO computation and counter-terms.

For the $\gamma\gamma$ states, according to our Refs. [29, 36], we have 4 polarization states that will be labeled as $\lambda_1\lambda_2 \in \{++, +-, -+, --\}$. The equation equivalent to Eq. (14), taking into account the effects of polarization, is

$$P_{IJ}^{\lambda_1\lambda_2} = \frac{1}{128\pi^2} \sqrt{\frac{4\pi}{2J+1}} \int d\Omega T_I^{\lambda_1\lambda_2}(s, \Omega) Y_{J,\Lambda}(\Omega), \quad \Lambda = \lambda_1 - \lambda_2. \quad (16)$$

Note that only $|I, M_I\rangle = |0, 0\rangle, |2, 0\rangle$ couples with 2- γ states, due to electric charge and angular momentum conservation. Indeed, parity conservation forbids 2- γ negative parity states coupling with $\{\omega^i\omega^j, hh\}$. Thus, let us introduce the positive parity state $(|+-\rangle + |-+\rangle)/\sqrt{2}$. This is the only one that couples with $J = 0$ states. Hence, we can define $P_{I0} \equiv (P_{I0}^{++} + P_{I0}^{--})$. For $J = 2$, the only non-vanishing contributions come from the (positive parity states) $|+-\rangle$ and $|-+\rangle$, so that $P_{I2} \equiv P_{I2}^{+-} = P_{I2}^{-+}$. The perturbative scattering amplitudes themselves be found in Ref. [29].

Finally, the $t^{\lambda_1}\bar{t}^{\lambda_2}$ states only couple with $|I = 0\rangle = \sum_i |\omega^i\omega^i\rangle/\sqrt{3}$ and hh states [37]. Even more, $|I = 0\rangle$ only couples with $|S = 1, S_Z = 0\rangle = (|+, +\rangle - |-, -\rangle)/\sqrt{2}$ $t\bar{t}$ state. The corresponding partial waves are

$$Q(\omega\omega \rightarrow t\bar{t}) = \frac{\sqrt{2}}{64\pi^2} \int d\Omega \mathcal{D}_{00}^0(\phi, \theta, -\phi) Q(\omega\omega \rightarrow t^+\bar{t}^+), \quad (17)$$

and a similar expression for $Q(hh \rightarrow t\bar{t})$.

4 Unitarization procedures

As pointed out in Fig. 1, our key assumption will be a strongly interacting EWSBS while couplings with $\gamma\gamma$ and $t\bar{t}$ states remain perturbative. The EWSBS $\{\omega^i\omega^j, hh\}$ scattering partial waves are unitarized by means of the same unitarization procedures (IAM, N/D, Improved-K matrix) that were exposed on our Refs. [28, 30].

We have two possibilities when unitarizing the EWSBS [30]. $\omega^i\omega^j$ states couple with hh if $IJ = 00, 20$ and some of these conditions are verified: $a^2 \neq b$, $d \neq 0$ or $e \neq 0$. Otherwise, the hh channel decouples. For instance, the single channel IAM is [30]

$$\tilde{A}_{IJ}^{\text{IAM}} = \frac{[A_{IJ}^{(0)}(s)]^2}{A_{IJ}^{(0)}(s) - A_{IJ}^{(1)}(s)} \quad (18)$$

Note that if coupling with hh channel happens, the coupled-channel versions of the unitarization procedures are required. Indeed, the strongly interacting regime could be triggered by $\omega\omega \rightarrow hh$ even if $\omega\omega \rightarrow \omega\omega$ is weak [28].

Couplings with $\gamma\gamma$ and $t\bar{t}$ are *unitarized* by means of a modification of the IAM and N/D procedures that uses as input the unitarized partial waves $\{\omega^i\omega^j, hh\} \rightarrow \{\omega^i\omega^j, hh\}$. This modification can be found in our Refs. [36, 37]. For the single-channel version of the IAM, it can be used

$$\tilde{P} = P^{(0)} \frac{\tilde{A}^{\text{IAM}}}{A^{(0)}}, \quad (19)$$

where \tilde{P} , in this case, is the unitarized $\omega\omega \rightarrow \{\gamma\gamma, t\bar{t}\}$ amplitude; $P^{(0)}$, the corresponding perturbative one; \tilde{A} , the unitarized $A(\omega\omega \rightarrow \omega\omega)$; and $A^{(0)}(s)$, the tree level matrix element $A(\omega\omega \rightarrow \omega\omega)$. This guarantees that the phase of \tilde{P} coincides with that of \tilde{A}^{IAM} . For other channels, see Refs. [28, 30].

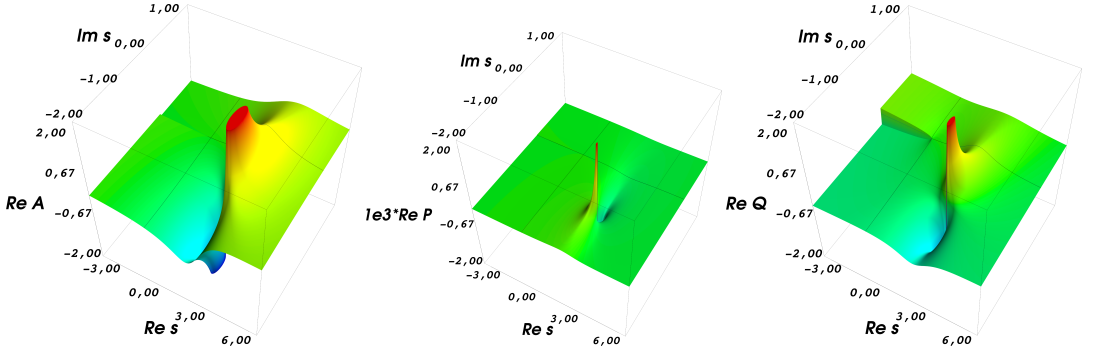


Figure 2. From left to right, second Riemann sheets of $\text{Im } A_{00}^{\text{IAM}}$, $10^3 \times \text{Im } P_{00}$ and $\text{Im } Q$. Parameters: $a^2 = b = 0.81^2$, $a4 = 4 \times 10^{-4}$, $g = 10^{-3}$. All the other NLO parameters set to 0.

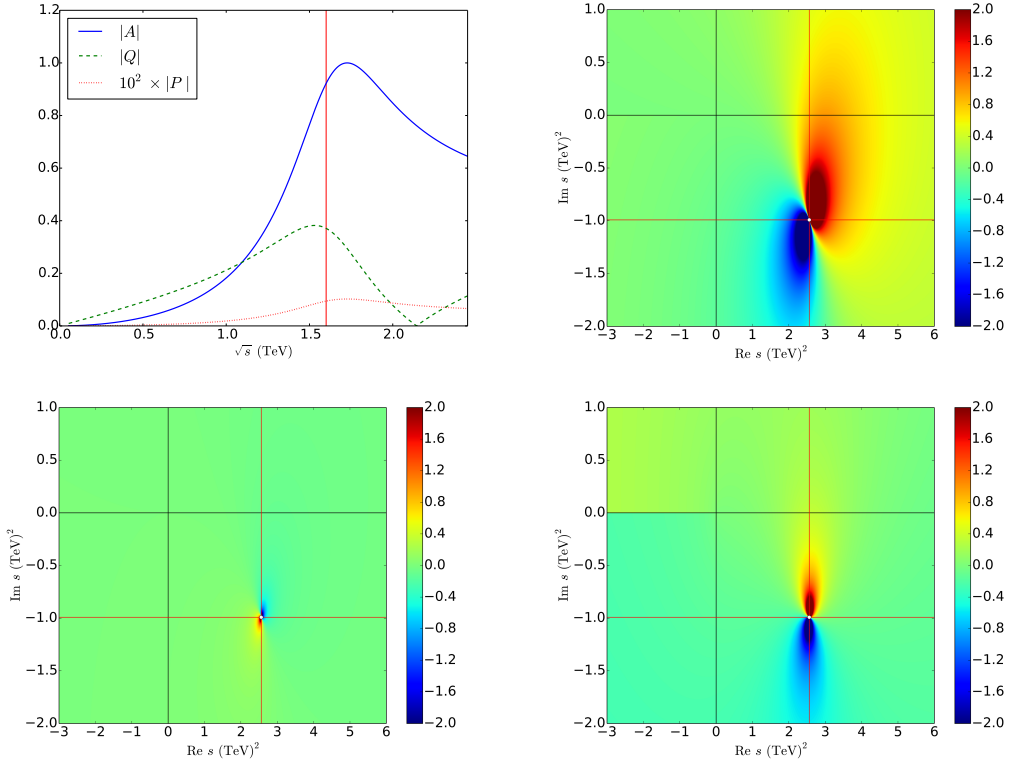


Figure 3. From top left, clockwise, $|A_{00}^{\text{IAM}}|$, $|Q|$ and $10^2 \times |P_{00}|^2$ over the physical region of s ; and second Riemann sheets of $\text{Im } A_{00}^{\text{IAM}}$, $\text{Im } Q$ and $\text{Im } P_{00}$. Same parameters of Fig. 2.

5 Numerical examples

On Figs. 2 and 3, a comparison between the $IJ = 00$ $A(\omega\omega \rightarrow \omega\omega)$, $Q(\omega\omega \rightarrow t\bar{t})$ and $P(\omega\omega \rightarrow \gamma\gamma)$ can be found. We have used $\mu = 3$ TeV (see matrix elements in Refs. [30, 36, 37]), $M_t = 175$ GeV, $a^2 = b = 0.81^2$, $a_4 = 4 \times 10^{-4}$, $g = 10^{-3}$ and all the other NLO counterterms set to 0. This corresponds to the unitarization of the single channel $\omega\omega \rightarrow \omega\omega$ plus weak couplings with $t\bar{t}$ and $\gamma\gamma$.

For computing the position of poles on the second Riemann sheet (resonances) in Fig. 3, we have used Cauchy's Theorem via numerical integration, following the procedure that was already explained in our Ref. [30].

Note the presence of a pole (broad resonance) on the second Riemann sheet in the same position in all the three channels (elastic scattering $\omega\omega \rightarrow \omega\omega$ and couplings with $t\bar{t}$ and $\gamma\gamma$). This was expected since the pole comes from the inner EWSBS (strongly coupled) dynamics, not from the physics of the (weak) coupling with $\gamma\gamma$ and $t\bar{t}$ states.

Poles in the second Riemann sheet, if there are any, should come from the unitarization procedures applied to the inner dynamics of the EWSBS (Eq. 18 and other expressions from our Ref. [30]). Couplings with $\gamma\gamma$ and $t\bar{t}$ do not give rise to poles in the second Riemann sheet, according to Eq. (19) (and similar expressions in Refs. [36, 37]), since perturbative partial waves (in particular, $P^{(0)}$ and $A^{(0)}$ in Eq. 19) are just polynomial in s/v^2 . Thus, the parameters associated with $\gamma\gamma$ (a_1, a_2, a_3, c_γ) and $t\bar{t}$ (c_1, c_2, g_t and g'_t) should control the $\gamma\gamma$ and $t\bar{t}$ coupling strength, but not the physics of the inner dynamics of the EWSBS. This is the picture which emerges from the chiral counting of Fig. 1. (However, if the parameters which control the couplings with $\gamma\gamma$ and $t\bar{t}$ states were unnaturally large, either M_t/v or α_{EW} would be non-perturbative and a full multi-channel unitarization involving $\{\gamma\gamma, t\bar{t}\}$ should be carried out [37].)

6 Conclusions

We have computed the NLO scattering processes $\{\omega\omega, hh\} \leftrightarrow \{\gamma\gamma, t\bar{t}\}$, and coupled them to a (hypothetical) strongly interacting EWSBS within the framework of HEFTs ($E < 4\pi v \sim 3$ TeV). The Equivalence Theorem [28–30] (which requires $E > M_h, M_W$) is used.

We have implemented all the relevant perturbative matrix elements and unitarization procedures for $\{\omega\omega, hh\} \leftrightarrow \{\omega\omega, hh, \gamma\gamma, t\bar{t}\}$ processes inside Fortran modules, following our Refs. [28–30, 36, 37]. Several unitarization procedures (based on dispersion relations) are used: IAM, N/D, Improve-K matrix and perturbative couplings with $\gamma\gamma$ and $t\bar{t}$ states.

Our formalism assumes that the inner EWSBS dynamics (in the W_L^\pm, Z_L and h sector) is stronger than their electromagnetic coupling with γS ($\sim \alpha_{EM}$) and the electroweak couplings with $t\bar{t}$ ($\sim M_t/v$). Hence, any new resonance (within the range $M_W, M_h < E < 3$ TeV) in $\{\omega\omega, hh\} \leftrightarrow \{\gamma\gamma, t\bar{t}\}$ channels would come from the inner EWSBS dynamics. This assumption could be broken for unnaturally large values of the NLO parameters related with $\gamma\gamma$ and $t\bar{t}$ couplings. Such a situation would require the usage of a full coupled channel unitarization, including strongly $\gamma\gamma$ or $t\bar{t}$ rescattering.

Even in the case that there are no new resonances, our scattering amplitudes can be a useful tool to parametrize separations from the SM in the regime $M_h, M_W < E < 3$ TeV. Currently, HEFTs themselves are being widely used in this way by CERN collaborations [23].

Finally, we are working within an expanded collaboration in the computation of simple estimates for collider cross-sections of typical resonances, and even for releasing a Monte Carlo module for the experimentalists.

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